

Baryon form factors

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Abstract

We calculate the form factors of the baryon octet in the framework of heavy baryon chiral perturbation theory. The calculated charge radius of the Σ^- is in agreement with recent measurements from CERN and Fermilab. We show that kaon loop effects can play a significant role in the neutron electric form factor. Furthermore, we derive generalized Caldi-Pagels relations between various charge radii which are free of chiral loop effects.

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1. Hadrons are composite objects which are characterized by certain sizes. The latter depend on the type of probe which is used to investigate the hadron structure. In the realm of strong QCD, the most precise information can be obtained using the electromagnetic (em) current. One measures the so-called electric and magnetic form factors, which in the Breit frame can be interpreted as the distribution of electric charge and magnetization within the hadron. The em form factors of the nucleon have been mapped out to high precision using electron scattering off protons and deuterons over many decades, for a comprehensive reference see e.g. [1]. The situation is very different for the other members of the ground state baryon octet, the hyperons. Until recently, for these only the magnetic moments had been determined. Last year, the first measurement of the Σ^- charge radius has been reported from experiments performed at CERN (WA89) [2] and FNAL (E781) [3] using electron scattering in inverse kinematics. More precisely, a very high energetic hyperon beam was scattered off atomic shell electrons and from the momentum dependence of the cross section the form factor could be extracted. For example, the SELEX collaboration [3] deduced from their data in the q^2 -region of -0.03 to -0.16 GeV² the mean square charge radius to be

$$\langle r^2 \rangle_{\Sigma^-} = 0.60 \pm 0.08_{\text{stat.}} \pm 0.08_{\text{syst.}} \text{ fm}^2. \quad (1)$$

Indeed, from naive quark model considerations one expects the hyperon size to be smaller than the one of the nucleon because the strange quark is more massive than the up and down quarks. A systematic investigation of the hyperon radii would allow one to study the deviations from flavor SU(3) as it can be done for the magnetic moments. Besides from these fundamental aspects, a knowledge of the hyperon form factors is also needed to interpret the kaon electroproduction data which are and will be obtained at ELSA (Bonn), MAMI-C (Mainz) and TJNAF. In fact, one can reverse the argument and try to use the processes $\gamma^* + N \rightarrow K + Y$, where γ^* (Y) denotes a virtual photon (hyperon), to determine the hyperon form factors (for a recent discussion, see ref. [4]). A related issue of current interest is the question of strangeness in the nucleon, which can either be inferred indirectly from kaon cloud contributions to the em nucleon form factors or more directly from a determination of the so-called strange nucleon form factors. It has been argued since a long time that in particular the neutron electric form factor might be particularly sensitive to strange quark effects, for an early reference see e.g. [5]. Furthermore, the strange magnetic ff was recently investigated using chiral perturbation theory and a leading one loop prediction free of adjustable parameters could be made [6]. Here, we use conventional three flavor baryon chiral perturbation theory in the heavy fermion formalism (called HBCHPT) to calculate the em form factors of the hyperons and the strangeness contributions in the nucleon. The momentum dependence of the ffs is governed by pion and kaon cloud contributions together with some counterterms which parametrize the effects of higher mass states. These effects can be calculated systematically and precisely. From previous studies of the nucleon in HBCHPT and extensions thereof [7], it is known that this approach is limited to low momentum transfer, say $-0.2 \leq q^2 \leq 0 \text{ GeV}^2$. Therefore, we will restrict our investigation to this region of momentum transfer. In the SU(3) study presented here, all appearing low-energy constants can be determined from the octet magnetic moments and the proton and neutron electric charge radii. Therefore, the momentum dependence of the em hyperon form factors can be predicted. Of course, many models have been used to calculate these form factors with wildly varying predictions, e.g. $0.3 \text{ fm}^2 \leq \langle r^2 \rangle_{\Sigma^-} \leq 1.2 \text{ fm}^2$, see e.g. fig.3 in [3]. The results obtained in HBCHPT allow one to further constrain these models.

2. The starting point of heavy baryon CHPT is an effective Lagrangian formulated in terms of the asymptotic fields, here the octet of Goldstone bosons and the ground state baryon octet B . The Lagrangian admits a low energy expansion of the form

$$\mathcal{L}_{\text{eff}} = \mathcal{L}_M + \mathcal{L}_{MB} = \mathcal{L}_M^{(2)} + \mathcal{L}_{MB}^{(1)} + \mathcal{L}_{MB}^{(2)} + \mathcal{L}_{MB}^{(3)} + \dots \quad (2)$$

where the subscript ' M ' (' MB ') denotes the meson (meson–baryon) sector and the superscript ' (i) ' the chiral dimension, which counts the power of external momenta and/or meson mass insertions. We call the corresponding small expansion parameter q . The Goldstone bosons are collected in the familiar $\text{SU}(3)$ –valued field $U = \exp\{i\phi/F_\phi\}$, with F_ϕ the octet decay constant. If not otherwise specified, we set $F_\phi = (F_\pi + F_K)/2 = 100 \text{ MeV}$. The effects due to $F_K \neq F_\pi$ are of higher order and can thus be neglected here. The baryons are given by a 3×3 matrix, transforming under $\text{SU}(3)_L \times \text{SU}(3)_R$ as usual matter fields, $B \rightarrow B' = KBK^\dagger$, with $K(U, L, R)$ the compensator field which is an element of the conserved subgroup $\text{SU}(3)_V$. We use the notation of [9]. The ellipsis stands for terms not needed here. Beyond leading order, the effective Lagrangian contains parameters not fixed by chiral symmetry, the so–called low–energy constants (LECs). These LECs must be be pinned down from data. In what follows, we will work to order q^3 in the chiral expansion. More precisely, we need the following terms from the second and third order meson–baryon Lagrangian (we omit finite terms contributing to the mass shift and alike):

$$\mathcal{L}_{MB}^{(2)} = -\frac{i}{4m} \langle \bar{B}[S^\mu, S^\nu][F_{\mu\nu}^+, B] \rangle - \frac{ib_{D/F}}{4m} \langle \bar{B}[S^\mu, S^\nu][F_{\mu\nu}^+, B]_\pm \rangle, \quad (3)$$

$$\begin{aligned} \mathcal{L}_{MB}^{(3)} &= \frac{b_{D/F}}{8m^2} \langle \bar{B}[[v^\mu D^\nu, F_{\mu\nu}^+], B]_\pm \rangle - \frac{d_{101/102}}{(4\pi F_\phi)^2} \langle \bar{B}[[v^\mu D^\nu, F_{\mu\nu}^+], B]_\mp \rangle \\ &+ \frac{id_{33/34}}{(4\pi F_\phi)^2} \left(\langle \bar{B}[\chi_+, [v \cdot D, B]]_\pm \rangle + \text{h.c.} \right) + \frac{id_{35}}{(4\pi F_\phi)^2} \left(\langle \bar{B}[v \cdot D, B] \rangle \langle \chi_+ \rangle + \text{h.c.} \right) \\ &+ \frac{id_{27}}{(4\pi F_\phi)^2} \langle \bar{B}[v \cdot D, [v \cdot D, [v \cdot D, B]]] \rangle, \end{aligned} \quad (4)$$

with S_μ the covariant spin–operator, $F_{\mu\nu}^+ = e(u^\dagger Q F_{\mu\nu} u + u Q F_{\mu\nu} u^\dagger) = 2e(\partial_\mu A_\nu - \partial_\nu A_\mu)Q + \mathcal{O}(\phi^2)$ and $\langle \dots \rangle$ denotes the trace in flavor space. Here, $Q = \text{diag}(2, -1, -1)/3$ is the quark charge matrix and $u = \sqrt{U}$. Throughout we work in the isospin limit $m_u = m_d$ and thus neglect the small $\Lambda - \Sigma^0$ mixing. The dimension two terms obviously give the magnetic coupling of the photon to the baryons. The corresponding LECs are determined with respect to the anomalous magnetic moments of the baryons. The third order LECs $d_{101,102}$ feature prominently in the electric radii whereas the last terms are only needed for renormalization and thus do not lead to observable effects.

3. In HBCHT, the natural frame to work with is the Breit–frame [8] and thus one defines the electromagnetic Sachs form factors $G_E(q^2)$ and $G_M(q^2)$. For a given baryon state from the octet, the matrix element of the electromagnetic current expressed in terms of the three light quark fields $q^T = (u, d, s)$, $V_\mu = \bar{q}Q\gamma_\mu q$, takes the form (to avoid unnecessary flavor indices, we work in the physical basis and consider mostly diagonal matrix elements, thus we can omit labelling the initial and final baryon states; the exception being the $\Lambda - \Sigma^0$ transition)

$$\langle B|V_\mu|B \rangle = \frac{1}{N_i N_f} \bar{u}(p') P_v^+ \left[G_E(q^2) v_\mu + \frac{1}{m} G_M(q^2) [S_\mu, S_\nu] q^\nu \right] P_v^+ u(p), \quad (5)$$

with

$$q_\mu = (p' - p)_\mu, \quad N = \sqrt{\frac{E + m}{2m}}, \quad (6)$$

and the P_v^+ are positive-velocity projection operators. For a more detailed discussion of this expression and the relation to the standard Dirac and Pauli form factors, see e.g. [7]. The chiral expansion of the form factors takes the form (here, G is a genuine symbol for any form factor)

$$G(q^2) = G^{\text{tree}}(q^2) + G^{\text{ct}}(q^2) + G^{\text{loop}}(q^2), \quad (7)$$

where we have split the tree graphs into the terms with fixed coefficients and the counterterms. For the magnetic form factors and the electric ones of the neutral octet baryons, the chiral expansion starts at second order. We thus expect to achieve the most precise description of the electric form factors for the charged particles. To third order, we have to deal with four non-trivial one loop graphs, two of the tadpole and two of the self-energy type, see e.g. ref. [7]. In the electric form factors, divergences appear. These are canceled by the appropriate β -functions of the operators 101 and 102 from table 1 of ref.[9], $\beta_{101} = -(9 + 25D^2 + 45F^2)/36$, $\beta_{102} = -5DF/2$. The general structure of all Sachs form factors is

$$G_E(q^2) = Q + \frac{1}{(4\pi F_\phi)^2} \left\{ \alpha_0 + \sum_{X=\pi,K} \alpha_X \ln \frac{M_X}{\lambda} + \sum_{X=\pi,K} \left[-\beta_X \left(2M_X^2 - \frac{5}{4}q^2 \right) \right. \right. \\ \left. \left. + \gamma_X \left(M_X^2 - \frac{1}{4}q^2 \right) \right] I_E^X(q^2) - (2Qd_{101}^r(\lambda) + \alpha_D d_{102}^r(\lambda))q^2 \right\} + \frac{q^2}{4m^2} (Qb_F + \alpha_D b^D), \quad (8)$$

$$G_M(q^2) = Q(1 + b_F) + \alpha_D b_D + \frac{m}{16\pi F_\phi^2} \sum_{X=\pi,K} \beta_X \left[M_X + \left(M_X^2 - \frac{1}{4}q^2 \right) I_M^X(q^2) \right], \quad (9)$$

with

$$I_E^X(q^2) = \frac{1}{3} \int_0^1 dx \ln \left(1 - x(1-x) \frac{q^2}{M_X^2} \right), \quad I_M^X(q^2) = \int_0^1 dx \frac{1}{\sqrt{M_X^2 - x(1-x)q^2}}, \quad (10)$$

and λ is the scale of dimensional regularization. Throughout, we set $\lambda = 1 \text{ GeV}$. The corresponding squared electric radii and magnetic slopes^{#4} are given by

$$r_E^2 = \frac{6}{Q + \delta_{Q0}} \frac{dG_E(q^2)}{dq^2} \Big|_{q^2=0} = \frac{1}{(4\pi F_\phi)^2} \left\{ \eta + 6\alpha_\pi \log \left(\frac{M_\pi}{\lambda} \right) + 6\alpha_K \log \left(\frac{M_K}{\lambda} \right) \right. \\ \left. - 12(Qd_{101}^r(\lambda) + \alpha_D d_{102}^r(\lambda)) \right\} + \frac{3}{2m^2} (Qb^F + \alpha_D b^D) \quad (11)$$

$$G'_M(0) = \frac{dG_M(q^2)}{dq^2} \Big|_{q^2=0} = -\frac{1}{96\pi F_\phi^2} \left\{ \beta_\pi \frac{m}{M_\pi} + \beta_K \frac{m}{M_K} \right\} \quad (12)$$

where $\eta = 6\alpha_0 + \frac{2}{3}\beta_\pi - \frac{1}{3}\gamma_\pi + \frac{2}{3}\beta_K - \frac{1}{3}\gamma_K$. Note that the normalization of the electric radii is such that one divides by the charge for a charged particle, i.e. by $G(0) = Q$, and by one for the

^{#4}We work with these slopes instead of the magnetic radii to avoid the uncertainty arising from the description of the magnetic moments to this order.

neutrals. The magnetic moments $\mu \equiv G_M(0)$ are

$$\mu = Q + \kappa = Q \left(1 + b^F\right) + \alpha_D b^D + \beta_\pi \frac{m M_\pi}{8\pi F_\phi^2} + \beta_K \frac{m M_K}{8\pi F_\phi^2}, \quad (13)$$

with κ the corresponding anomalous magnetic moment. The pertinent coefficients are shown in the table 1.

	Q	α_0	α_π	α_K
p	1	$-\frac{1}{12} - \frac{85}{108}D^2 - \frac{17}{18}DF - \frac{17}{12}F^2$	$-\frac{1}{6} - \frac{5}{6}(D+F)^2$	$-\frac{1}{3} - \frac{5}{3}\left(\frac{D^2}{3} + F^2\right)$
n	0	$\frac{17}{9}DF$	$\frac{1}{6} + \frac{5}{6}(D+F)^2$	$-\frac{1}{6} - \frac{5}{6}(D-F)^2$
Σ^+	1	$-\frac{1}{12} - \frac{85}{108}D^2 - \frac{17}{18}DF - \frac{17}{12}F^2$	$-\frac{1}{3} - \frac{5}{3}\left(\frac{D^2}{3} + F^2\right)$	$-\frac{1}{6} - \frac{5}{6}(D+F)^2$
Σ^-	-1	$\frac{1}{12} + \frac{85}{108}D^2 - \frac{17}{18}DF + \frac{17}{12}F^2$	$\frac{1}{3} + \frac{5}{3}\left(\frac{D^2}{3} + F^2\right)$	$\frac{1}{6} + \frac{5}{6}(D-F)^2$
Σ^0	0	$-\frac{17}{18}DF$	0	$-\frac{5}{3}DF$
Λ	0	$\frac{17}{18}DF$	0	$\frac{5}{3}DF$
$\Lambda\Sigma^0$	0	$-\frac{17}{6\sqrt{3}}DF$	$-\frac{10}{3\sqrt{3}}DF$	$-\frac{5}{3\sqrt{3}}DF$
Ξ^0	0	$\frac{17}{9}DF$	$-\frac{1}{6} - \frac{5}{6}(D-F)^2$	$\frac{1}{6} + \frac{5}{6}(D+F)^2$
Ξ^-	-1	$\frac{1}{12} + \frac{85}{108}D^2 - \frac{17}{18}DF + \frac{17}{12}F^2$	$\frac{1}{6} + \frac{5}{6}(D-F)^2$	$\frac{1}{3} + \frac{5}{3}\left(\frac{D^2}{3} + F^2\right)$

	α_D	β_π	γ_π	β_K	γ_K	η
p	$\frac{1}{3}$	$-(D+F)^2$	1	$-2\left(\frac{D^2}{3} + F^2\right)$	2	$-\frac{3}{2} - \frac{35}{6}D^2 - 7DF - \frac{21}{2}F^2$
n	$-\frac{2}{3}$	$(D+F)^2$	-1	$-(D-F)^2$	1	$14DF$
Σ^+	$\frac{1}{3}$	$-2\left(\frac{D^2}{3} + F^2\right)$	2	$-(D+F)^2$	1	$-\frac{3}{2} - \frac{35}{6}D^2 - 7DF - \frac{21}{2}F^2$
Σ^-	$-\frac{1}{3}$	$2\left(\frac{D^2}{3} + F^2\right)$	-2	$(D-F)^2$	-1	$\frac{3}{2} + \frac{35}{6}D^2 - 7DF + \frac{21}{2}F^2$
Σ^0	$\frac{1}{3}$	0	0	$-2DF$	0	$-7DF$
Λ	$-\frac{1}{3}$	0	0	$2DF$	0	$7DF$
$\Lambda\Sigma^0$	$\frac{1}{\sqrt{3}}$	$-\frac{4}{\sqrt{3}}DF$	0	$-\frac{2}{\sqrt{3}}DF$	0	$-7\sqrt{3}DF$
Ξ^0	$-\frac{2}{3}$	$-(D-F)^2$	1	$(D+F)^2$	-1	$14DF$
Ξ^-	$\frac{1}{3}$	$(D-F)^2$	-1	$2\left(\frac{D^2}{3} + F^2\right)$	-2	$\frac{3}{2} + \frac{35}{6}D^2 - 7DF + \frac{21}{2}F^2$

Table 1: Table of coefficients for the various octet states.

We remark that the last term in eq.(8) can be expressed in terms of the anomalous magnetic moment of the baryon considered by use of eq.(13), i.e. it is the well-known Foldy term. In the limit $M_K \rightarrow \infty$, we recover the SU(2) results for the form factors as demanded by decoupling. It is worth to notice that we can derive three relations for the magnetic moments, electric and magnetic charge radii which are free of chiral loop effects (these have been found originally for the magnetic moments only in ref.[10]). Denoting by O the observables r_E^2 , r_M^2 and μ , these generalized Caldi-Pagels relations take the form

$$\frac{1}{2}O(\Sigma^+ + \Sigma^-) = O(\Sigma^0) = -O(\Lambda), \quad (14)$$

$$O(\Xi^0 + \Xi^- + p + n) = 2O(\Lambda) , \quad (15)$$

$$O(\Xi^0 + n + \sqrt{3}\Lambda\Sigma^0) = O(\Lambda) . \quad (16)$$

The first of these is nothing but isospin symmetry. Note that for the magnetic moments, these relations work astonishingly well, within a few percent.

4. Before presenting results, we must fix parameters. Throughout, we use $D = 3/4$, $F = 1/2$ and $F_\phi = 100$ MeV. The other parameters are varied as follows

$$\begin{aligned} \text{Set 1} & : b_D = 3.92, b_F = 2.92, d_{101}^r(1 \text{ GeV}) = -1.06, d_{102}^r(1 \text{ GeV}) = 1.70, m = 0.94 \text{ GeV} , \\ \text{Set 2} & : b_D = 5.17, b_F = 2.76, d_{101}^r(1 \text{ GeV}) = -1.06, d_{102}^r(1 \text{ GeV}) = 1.70, m = 0.94 \text{ GeV} , \\ \text{Set 3} & : b_D = 5.81, b_F = 3.22, d_{101}^r(1 \text{ GeV}) = -1.27, d_{102}^r(1 \text{ GeV}) = 1.68, m = 1.15 \text{ GeV} . \end{aligned} \quad (17)$$

Here, the first set is chosen such that the proton and neutron electric radii and magnetic moments are exactly reproduced. In set 2, an overall best fit to the octet magnetic moments is achieved without readjusting the LECs $d_{101,102}$. In the third set, the average octet mass is used while $b_{D,F}$ are obtained from the octet magnetic moments and $d_{101,102}$ are adjusted to reproduce $r_{E,p}^2$ and $r_{E,n}^2$. We note that for all sets the LECs $d_{101,102}^r(\lambda = 1 \text{ GeV})$ are of natural size. While sets 1,3 lead to the same octet charge radii, the sets 1,2 obviously give the same magnetic slopes. The resulting charge radii and magnetic slopes are collected in table 2.

B	$r_E^2 [\text{fm}^2]$			$G'_M(0) [\text{fm}^2]$		
	Set 1	Set 2	Exp.	Set 1	Set 3	Exp.
p	0.735	0.717	0.735 [1]	0.157	0.192	0.325 [1]
n	-0.113	-0.168	-0.113±0.004 [11]	-0.134	-0.164	0.252 [1]
Σ^+	0.642	0.659	—	0.114	0.140	—
Σ^-	0.803	0.765	0.60±0.08±0.08 [3] 0.91±0.32±0.40 [2]	-0.077	-0.095	—
Σ^0	-0.081	-0.053	—	0.018	0.023	—
Λ	0.081	0.053	—	-0.018	-0.023	—
Ξ^0	0.202	0.147	—	-0.033	-0.040	—
Ξ^-	0.645	0.608	—	-0.027	-0.033	—
$\Lambda\Sigma^0$	-0.005	0.004	—	0.086	0.150	—

Table 2: Electric charge radii and magnetic slopes for the octet baryons for the parameter sets described in the text.

Let us first discuss the implications for the nucleon sector (using the results obtained with set 1). The effect of the kaon cloud on the proton charge form factor is negligible, whereas the description of the neutron charge form factor is obviously improved, as shown in fig.1. This is

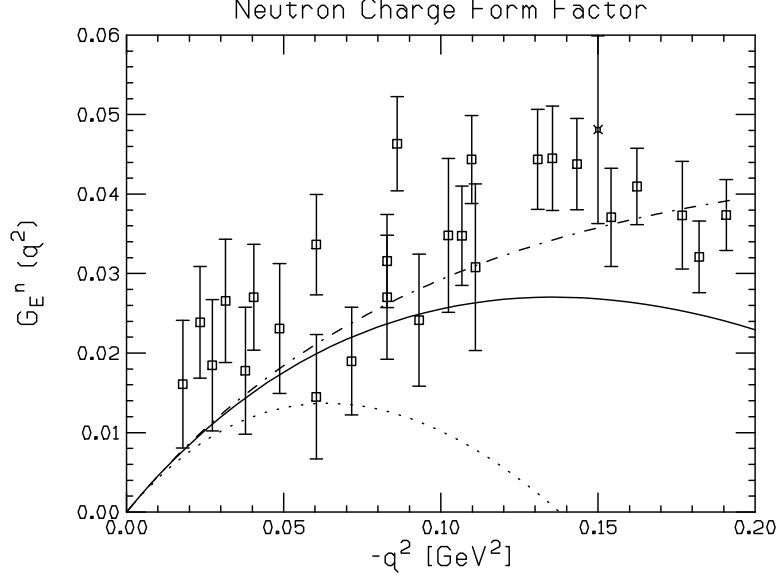


Figure 1: The neutron electric form factor. The solid and the dotted lines are the third order SU(3) and SU(2) HBCPT predictions, respectively. The model-dependent data are taken from [12] and the dot-dashed line is the dispersion-theoretical fit of [1]. The filled cross indicates the result of the Mainz $D(\vec{e}, e'\vec{n})$ experiment [13].

mostly due to the fact that in SU(3), the kaon cloud induces an additional momentum-dependence of the isoscalar electric and magnetic form factors. In contrast, in SU(2) the isoscalar magnetic ff is simply a constant to third order (as it was already pointed out in ref.[6]). Most likely, this sizeable kaon cloud effect will be reduced when higher orders are calculated. Also, the strangeness effects disimprove the description of the magnetic form factors compared to the SU(2) results. As already stressed, the magnetic coupling of the photon only starts at second order and one thus should go to fourth order for a precise calculation of the magnetic form factors (this is explicitly demonstrated for the magnetic moments in [9]). It is also known from the SU(2) calculation in ref. [7] that the contribution from intermediate decuplet states brings the magnetic radii closer to their empirical values. In HBCPT, such effects only start to appear at fourth order.

The charge radii and magnetic slopes of the hyperons can also be found in table 2. The resulting numbers for the electric charge radii are fairly stable under parameter variations for the charged hyperons. We stress that the one for the Σ^- agrees within the quoted uncertainty with the measured values obtained at CERN and FNAL. The corresponding electric and magnetic form factors of all hyperons are collected in figs. 2,3 for parameter set 3. It is worth to point out that the electric ffs of the neutral hyperons do not show the bending as the neutron electric ff does. This can be traced back to the suppression of the pion cloud component for these particular baryons. From the arguments given above, we expect the magnetic slopes to be increased in magnitude when the calculation is done to next order. The quality of the prediction for the magnetic form factors is expected to be similar to the one for the magnetic moments.

5. We have investigated the electromagnetic form factors of the ground state hyperon octet to

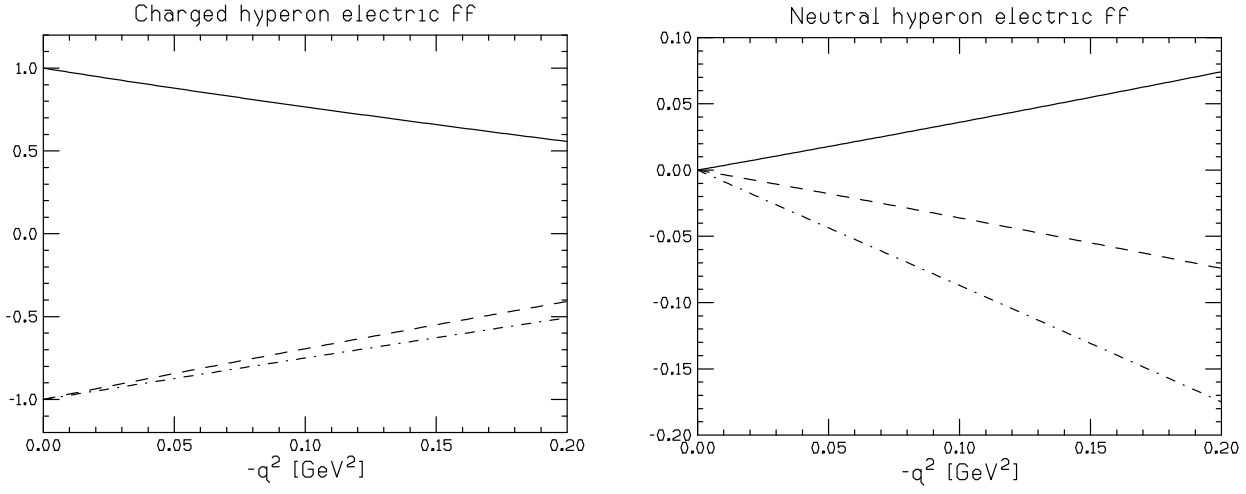


Figure 2: Electric hyperon form factors. Left panel: Charged particles, Σ^+ (solid line), Σ^- (dashed line), Ξ^- (dot-dashed line). Right panel: Neutral particles, Σ^0 (solid line), Λ (dashed line), Ξ^0 (dot-dashed line).

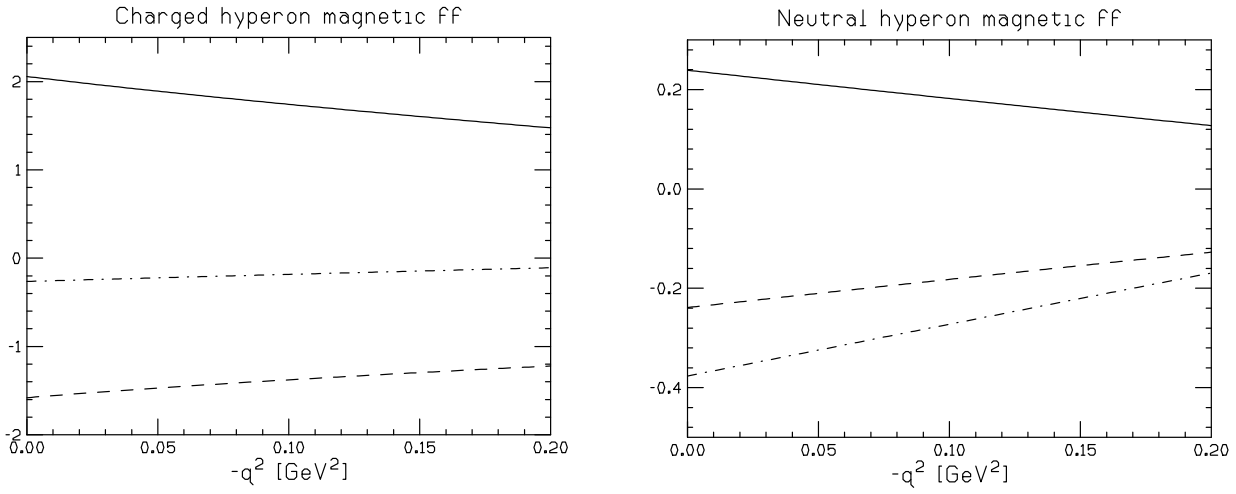


Figure 3: Magnetic hyperon form factors. Left panel: Charged particles, Σ^+ (solid line), Σ^- (dashed line), Ξ^- (dot-dashed line). Right panel: Neutral particles, Σ^0 (solid line), Λ (dashed line), Ξ^0 (dot-dashed line).

leading one loop order in heavy baryon chiral perturbation theory. To this order, one has four low-energy constants. Two of these can be fixed from the magnetic moments and the other two from the electric radii of the proton and the neutron. With that, the momentum-dependence of the hyperon form factors is given parameter-free. The prediction for the Σ^- radius agrees with the recently obtained value at CERN and FNAL using high energetic hyperon beams. We have also derived the so-called generalized Caldi-Pagels relations between various charge radii. These are free of chiral loop effects and might eventually lead to constraints on some neutral hyperon radii. We also have investigated the strangeness contribution to the nucleon electromagnetic form factors and showed that strangeness can contribute significantly to the neutron charge form factor. A next order calculation is called for to substantiate such a claim.

Acknowledgements

We thank Hans-Werner Hammer for a careful reading of the manuscript. One of us (UGM) thanks the Institute for Nuclear Theory at the University of Washington for its hospitality and the DOE for partial support during the completion of this work.

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